



A canonical $D_s(2317)$?

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Abstract

It is shown that quark mass dependence induced by one-loop corrections to the Breit–Fermi spin-dependent one-gluon-exchange potential permits an accurate determination of heavy–light meson masses. Thus the $D_s(2317)$ is a canonical $c\bar{s}$ meson in this scenario. The multiplet splitting relationship of chiral doublet models, $M(1^+) - M(1^-) = M(0^+) - M(0^-)$, holds to good accuracy in the D and D_s systems, but is accidental. Radiative transitions and bottom flavoured meson masses are discussed.

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1. Introduction

BaBar’s discovery of the $D_s(2317)$ state [1] generated strong interest in heavy meson spectroscopy, chiefly due to its surprisingly low mass with respect to expectations. These expectations are based on quark models or lattice gauge theory. Unfortunately, at present large lattice systematic errors do not allow a determination of the D_s mass with a precision better than several hundred MeV. And, although quark models appear to be exceptionally accurate in describing charmonia, they are less constrained by experiment and on a weaker theoretical footing in the open charm sector. It is therefore imperative to examine reasonable alternative descriptions of the open charm sector.

The $D_s(2317)$ was produced in e^+e^- scattering and discovered in the isospin violating $D_s\pi$ decay mode in $K\bar{K}\pi\pi$ and $K\bar{K}\pi\pi\pi$ mass distributions. Its width is less than 10 MeV and it is likely that the quantum numbers are $J^P = 0^+$ [2]. Finally, if the $D_s\pi^0$ mode dominates the width of the $D_s(2317)$ then the measured product of branching ratios [3]

$$Br(B^0 \rightarrow D_s(2317)K) \cdot Br(D_s(2317) \rightarrow D_s\pi^0) = (4.4 \pm 0.8 \pm 1.1) \times 10^{-5} \quad (1)$$

implies that $Br(B \rightarrow D_s(2317)K) \approx Br(B \rightarrow D_s K)$, consistent with the $D_s(2317)$ being a canonical $0^+ c\bar{s}$ meson.

In view of this, Cahn and Jackson have examined the feasibility of describing the masses and decay widths of the low lying D and D_s states within the constituent quark model [4]. They assume a standard spin-dependent structure for the quark–antiquark interaction (see below) and allow general vector and scalar potentials. Their conclusion is that it is very difficult to describe the data in this scenario.

Indeed, the $D_s(2317)$ lies some 160 MeV below most model predictions (see Ref. [2] for a summary), leading to speculation that the state could be a DK molecule [5] or a tetraquark [6]. Such speculation is supported by the isospin violating discovery mode of the $D_s(2317)$ and the proximity of the S-wave DK threshold at 2358–2367 MeV. Other studies have been made with QCD sum rules [7], using heavy quark symmetry to examine decay models [8], or in unitarised chiral models [9].

Although these proposals have several attractive features, it is important to exhaust possible canonical $c\bar{s}$ descriptions of the $D_s(2317)$ before resorting to more exotic models. Here we propose a simple modification to the standard vector Coulomb + scalar

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Table 1
Low lying D and D_s spectra (MeV)

J^P	0^-	1^-	0^+	1^+	1^+	2^+
D	1869.3 ± 0.5	2010.0 ± 0.5	a	b	2422.2 ± 1.8	2459 ± 4
D_s	1968.5 ± 0.6	2112.4 ± 0.7	2317.4 ± 0.9	2459.3 ± 1.3	2535.35 ± 0.34	2572.4 ± 1.5

linear quark potential model that maintains good agreement with the charmonium spectrum and agrees remarkably well with the D and D_s spectra. Possible experimental tests of this scenario are discussed.

2. A quark model of open charm states

The spectra we seek to explain are summarised in Table 1. Unfortunately, the masses of the D_0 (labelled a) and D_1 (labelled b) are poorly determined. Belle have observed [10] the D_0 in B decays and claim a mass of $2308 \pm 17 \pm 32$ MeV with a width of $\Gamma = 276 \pm 21 \pm 18 \pm 60$, while FOCUS [11] find $2407 \pm 21 \pm 35$ MeV with a width $\Gamma = 240 \pm 55 \pm 59$. While some authors choose to average these values, we regard them as incompatible and consider the cases separately below. Finally, there is an older mass determination from Belle [12] of $2290 \pm 22 \pm 20$ MeV with a width of $\Gamma = 305 \pm 30 \pm 25$. The D'_1 has been seen in B decays to $D\pi\pi$ and $D^*\pi\pi$ by Belle [13]. A Breit–Wigner fit yields a mass of $2427 \pm 26 \pm 20 \pm 15$ MeV and a width of $384^{+107}_{-90} \pm 24 \pm 70$ MeV. Alternatively, a preliminary report from CLEO [14] cites a mass of $2461^{+41}_{-34} \pm 10 \pm 32$ MeV and a width of $290^{+101}_{-79} \pm 26 \pm 36$ MeV. Finally, FOCUS [15] obtain a lower neutral D'_1 mass of $2407 \pm 21 \pm 35$ MeV. Other masses in Table 1 are obtained from the PDG compilation [16].

In addition to the unexpectedly low mass of the $D_s(2317)$, the $D_s(2460)$ is also somewhat below predictions assuming it is the D_{s1} (Godfrey and Isgur, for example, predict a mass of 2530 MeV [17]). It is possible that an analogous situation holds in the D spectrum, depending on the mass of the D_0 . The quark model explanation of these states rests on P-wave mass splittings induced by spin-dependent interactions. A common model of spin-dependence is based on the Breit–Fermi reduction of the one-gluon-exchange interaction supplemented with the spin-dependence due to a scalar current confinement interaction. The general form of this potential has been computed by Eichten and Feinberg [18] at tree level using Wilson loop methodology. The result is parameterised in terms of four nonperturbative matrix elements, V_i , which can be determined by electric and magnetic field insertions on quark lines in the Wilson loop. Subsequently, Gupta and Radford [19] performed a one-loop computation of the heavy quark interaction and showed that a fifth interaction, V_5 is present in the case of unequal quark masses. The net result is a quark–antiquark interaction that can be written as:

$$V_{q\bar{q}} = V_{\text{conf}} + V_{\text{SD}} \quad (2)$$

where V_{conf} is the standard Coulomb + linear scalar form:

$$V_{\text{conf}}(r) = -\frac{4}{3} \frac{\alpha_s}{r} + br \quad (3)$$

and

$$\begin{aligned} V_{\text{SD}}(r) = & \left(\frac{\sigma_q}{4m_q^2} + \frac{\sigma_{\bar{q}}}{4m_{\bar{q}}^2} \right) \cdot \mathbf{L} \left(\frac{1}{r} \frac{dV_{\text{conf}}}{dr} + \frac{2}{r} \frac{dV_1}{dr} \right) + \left(\frac{\sigma_{\bar{q}} + \sigma_q}{2m_q m_{\bar{q}}} \right) \cdot \mathbf{L} \left(\frac{1}{r} \frac{dV_2}{dr} \right) \\ & + \frac{1}{12m_q m_{\bar{q}}} (3\sigma_q \cdot \hat{\mathbf{r}} \sigma_{\bar{q}} \cdot \hat{\mathbf{r}} - \sigma_q \cdot \sigma_{\bar{q}}) V_3 + \frac{1}{12m_q m_{\bar{q}}} \sigma_q \cdot \sigma_{\bar{q}} V_4 \\ & + \frac{1}{2} \left[\left(\frac{\sigma_q}{m_q^2} - \frac{\sigma_{\bar{q}}}{m_{\bar{q}}^2} \right) \cdot \mathbf{L} + \left(\frac{\sigma_q - \sigma_{\bar{q}}}{m_q m_{\bar{q}}} \right) \cdot \mathbf{L} \right] V_5. \end{aligned} \quad (4)$$

Here $\mathbf{L} = \mathbf{L}_q = -\mathbf{L}_{\bar{q}}$, $r = |\mathbf{r}| = |\mathbf{r}_q - \mathbf{r}_{\bar{q}}|$ is the $\bar{Q}Q$ separation and the $V_i = V_i(m_q, m_{\bar{q}}; r)$ are the Wilson loop matrix elements discussed above.

The first four V_i are order α_s in perturbation theory, while V_5 is order α_s^2 ; for this reason V_5 has been largely ignored by quark modellers. The exceptions known to us are Ref. [20], which examines S-wave masses for a variety of heavy–light mesons in a model very similar to that presented here, and the second of Ref. [19], which does not consider scalar confinement contributions to the spin-dependent interaction. More recently, Cahn and Jackson [4] only consider $V_1 - V_4$ in an analysis of the D_s system. In practise this is acceptable (as we show below) *except in the case of unequal quark masses*, where the additional spin–orbit interaction can play an important role.

Here we propose to take the spin-dependence of Eq. (4) seriously and examine its effect on low-lying heavy–light mesons. Our model can be described in terms of vector and scalar kernels defined by

$$V_{\text{conf}} = V + S \quad (5)$$

where $V = -4\alpha_s/3r$ is the vector kernel and $S = br$ is the scalar kernel, and by the order α_s^2 contributions to the V_i , denoted by δV_i . Expressions for the matrix elements of the spin-dependent interaction are then

$$\begin{aligned} V_1 &= -S + \delta V_1, \\ V_2 &= V + \delta V_2, \\ V_3 &= V'/r - V'' + \delta V_3, \\ V_4 &= 2\nabla^2 V + \delta V_4, \\ V_5 &= \delta V_5. \end{aligned} \quad (6)$$

Explicitly,

$$\begin{aligned} V_1(m_q, m_{\bar{q}}, r) &= -br - C_F \frac{1}{2r} \frac{\alpha_s^2}{\pi} (C_F - C_A (\ln[(m_q m_{\bar{q}})^{1/2} r] + \gamma_E)), \\ V_2(m_q, m_{\bar{q}}, r) &= -\frac{1}{r} C_F \alpha_s \left[1 + \frac{\alpha_s}{\pi} \left[\frac{b_0}{2} [\ln(\mu r) + \gamma_E] + \frac{5}{12} b_0 - \frac{2}{3} C_A + \frac{1}{2} (C_F - C_A (\ln[(m_q m_{\bar{q}})^{1/2} r] + \gamma_E)) \right] \right], \\ V_3(m_q, m_{\bar{q}}, r) &= \frac{3}{r^3} C_F \alpha_s \left[1 + \frac{\alpha_s}{\pi} \left[\frac{b_0}{2} \left[\ln(\mu r) + \gamma_E - \frac{4}{3} \right] + \frac{5}{12} b_0 - \frac{2}{3} C_A \right. \right. \\ &\quad \left. \left. + \frac{1}{2} \left(C_A + 2C_F - 2C_A (\ln[(m_q m_{\bar{q}})^{1/2} r] + \gamma_E - \frac{4}{3}) \right) \right] \right], \\ V_4(m_q, m_{\bar{q}}, r) &= \frac{32\alpha_s \sigma^3 e^{-\sigma^2 r^2}}{3\sqrt{\pi}}, \\ V_5(m_q, m_{\bar{q}}, r) &= \frac{1}{4r^3} C_F C_A \frac{\alpha_s^2}{\pi} \ln \frac{m_{\bar{q}}}{m_q} \end{aligned} \quad (7)$$

where $C_F = 4/3$, $C_A = 3$, $b_0 = 9$, $\gamma_E = 0.5772$, and the scale μ has been set to 1 GeV.

The hyperfine interaction (proportional to V_4) contains a delta function in configuration space and is normally ‘smeared’ to make it nonperturbatively tractable. This introduces a new parameter that largely subsumes corrections to the hyperfine interaction such as δV_4 . For this reason we choose not to include δV_4 in the model definition of Eq. (7). Corrections to the remaining terms are included because they retain their perturbative forms. In the following, the hyperfine interaction has been treated nonperturbatively and the remaining spin-dependent terms are evaluated in leading-order perturbation theory.

We have confirmed that the additional features do not ruin previous agreement with the charmonium spectrum. For example, Ref. [21] obtains very good agreement with experiment for parameters $m_c = 1.4794$ GeV, $\alpha_s = 0.5461$, $b = 0.1425$ GeV², and $\sigma = 1.0946$ GeV. Employing the model of Eq. (7) worsens the agreement with experiment, but the original good fit is recovered upon slightly modifying parameters (the refit parameters are $m_c = 1.57$ GeV, $\alpha_s = 0.52$, $b = 0.15$ GeV², and $\sigma = 1.3$ GeV).

The low lying $c\bar{s}$ and $c\bar{u}$ states are fit reasonably well with the parameters labelled ‘avg’ in Table 2. Predicted masses are given in Table 3. Parameters labelled ‘low’ in Table 2 fit the D mesons very well, whereas those labelled ‘high’ fit the known D_s mesons well. It is thus reassuring that these parameter sets are reasonably similar to each other and to the refit charmonium parameters. (Note that constant shifts in each flavour sector are determined by the relevant pseudoscalar masses.)

The predicted D_{s0} mass is 2341 MeV, 140 MeV lower than the prediction of Godfrey and Isgur and only 24 MeV higher than experiment. We remark that the best fit to the D spectrum predicts a mass of 2287 MeV for the D_0 meson, in good agreement with the preliminary Belle measurement of 2290 MeV, 21 MeV below the current Belle mass, and in disagreement with the FOCUS mass of 2407 MeV.

Table 2
Model parameters

Model	α_s	b (GeV ²)	σ (GeV)	m_c (GeV)	C (GeV)
low	0.46	0.145	1.20	1.40	−0.298
avg	0.50	0.140	1.17	1.43	−0.275
high	0.53	0.135	1.13	1.45	−0.254

Table 3
Predicted low lying charm meson spectra (GeV)

Flavour	0^-	1^-	0^+	1^+	1^+	2^+
D	1.869	2.017	2.260	2.406	2.445	2.493
D_s	1.968	2.105	2.341	2.475	2.514	2.563

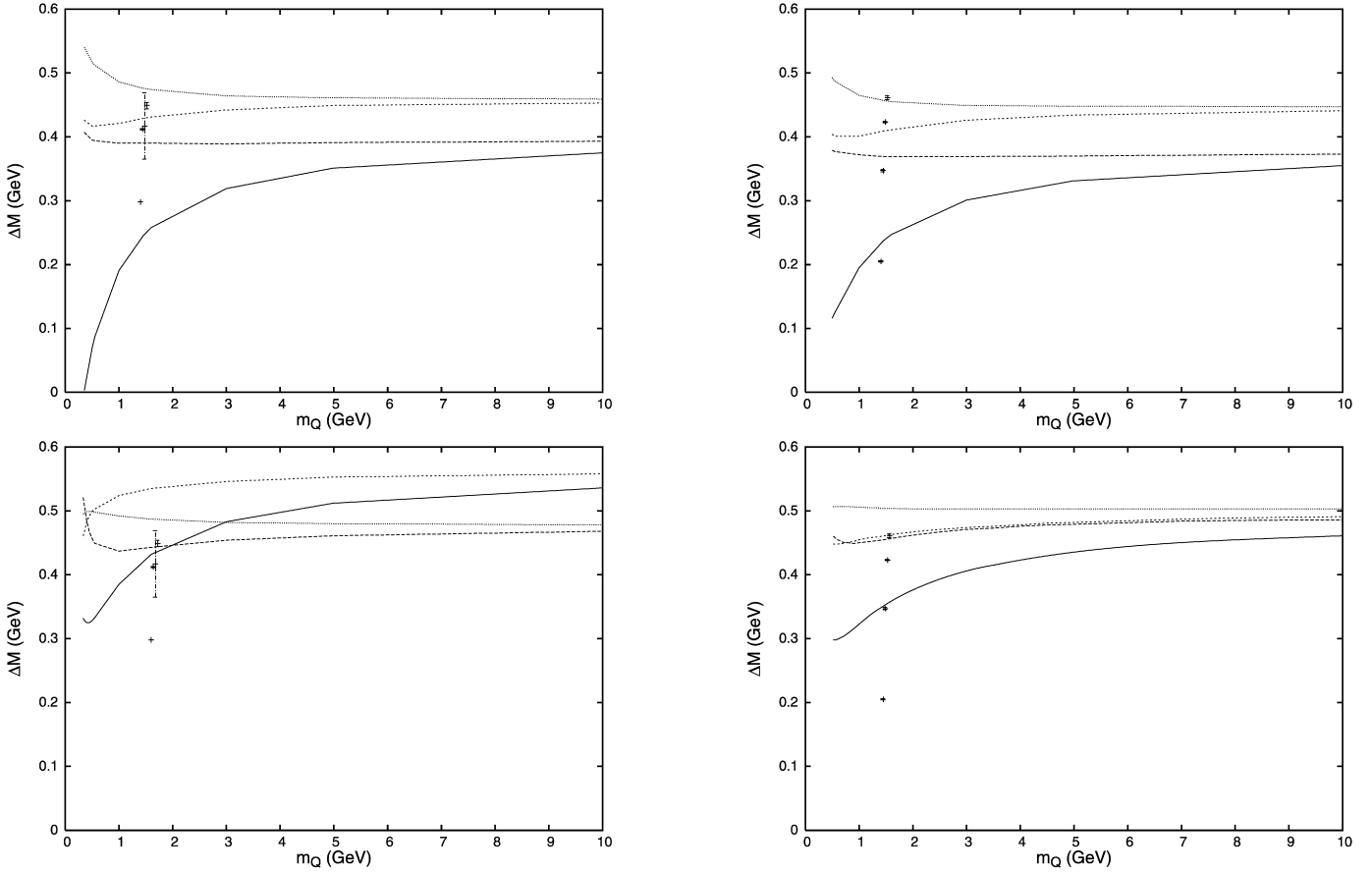


Fig. 1. $M(\text{P-wave}) - M(\text{vector})$ as a function of the heavy quark mass. D system (left); D_s system (right); one-loop model (top); traditional model (bottom). D_2 (dotted line); D'_1 (thin dashed line); D_1 (dashed line); D_0 (solid line). The points are experimental mass splittings.

The average error in the predicted P-wave masses is less than 1%. It thus appears likely that a simple modification to the spin-dependent quark interaction is capable of describing heavy–light mesons with reasonable accuracy.

We examine the new model in more detail by computing P-wave meson masses (with respect to the ground state vector) as a function of the heavy quark mass. Results for $Q\bar{u}$ and $Q\bar{s}$ systems are displayed in Fig. 1. All panels indicate that the approach to the heavy quark limit is very slow. In the case of the traditional $Q\bar{u}$ system the heavy quark doublets are inverted (with the $j_q = 1/2$ doublet higher than the $j_q = 3/2$), in disagreement with experiment. Alternatively, the one-loop model displays the expected heavy quark behaviour. Furthermore, the predicted mass splittings at the charm quark scale are near experiment for D masses (points on the panels). A similar situation holds for the D_s system (right panels), except in this case it is the D_s and D'_s that do not agree with the traditional model predictions.

Although the reliability of the model is suspect in the case of light Q masses, it is intriguing that the one-loop model scalar–vector mass difference gets small in this limit. Thus it is possible that the enigmatic a_0 and f_0 mesons may simply be $q\bar{q}$ states.

Finally, one obtains $M(h_c) > M(\chi_{c1})$ in one-loop and traditional models, in agreement with experiment. However, experimentally $M(f_1) - M(h_1) \approx 100$ MeV and $M(a_1) - M(b_1) \approx 0$ MeV, indicating that the 3P_1 state is heavier than (or nearly degenerate with) the 1P_1 light meson state. Thus the sign of the combination of tensor and spin–orbit terms that drives this splitting must change when going from charm quark to light quark masses. This change is approximately correctly reproduced in the traditional model (lower left panel of Fig. 1). The one-loop model does not reproduce the desired cross over, although it does come close, and manipulating model parameters can probably reproduce this behaviour. We do not pursue this here since the focus is on heavy–light mesons.

3. Mixing angles and radiative decays

The lack of charge conjugation symmetry implies that two nearby low lying axial vector states exist (generically denoted as D_1 and D'_1 in the following). The mixing angle between these states can be computed and compared to experiment (with the help of additional model assumptions). We define the mixing angle via the relations:

$$\begin{aligned} |D_1\rangle &= +\cos(\phi)|^1P_1\rangle + \sin(\phi)|^3P_1\rangle, \\ |D'_1\rangle &= -\sin(\phi)|^1P_1\rangle + \cos(\phi)|^3P_1\rangle. \end{aligned} \quad (8)$$

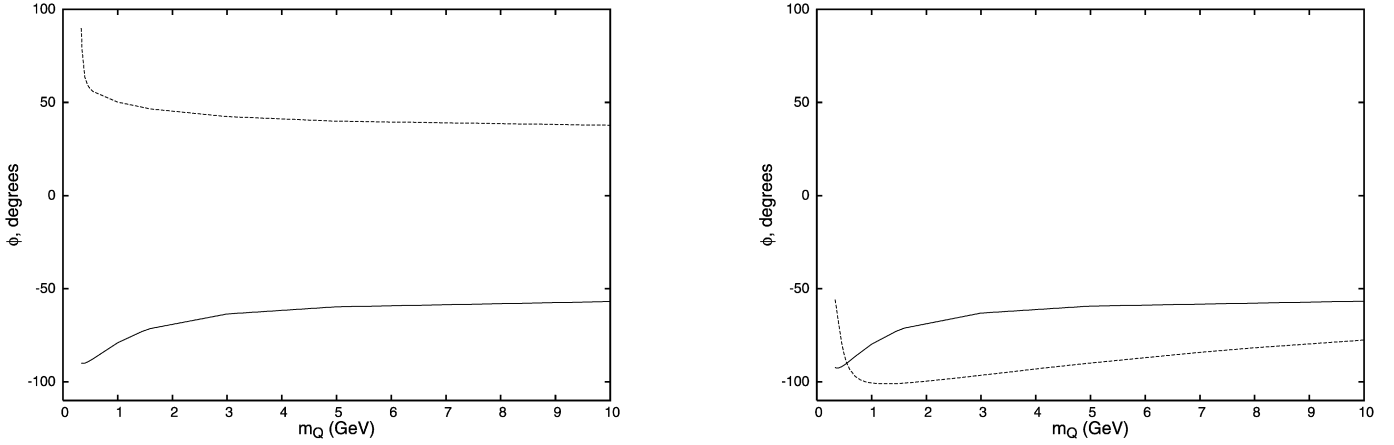


Fig. 2. D_1 (left) and D_{s1} (right) mixing angles vs. heavy quark mass. The traditional model is given by the dashed line; the extended model is the solid line.

In the following, we choose to define the D'_1 as the heavier axial state in the heavy quark limit. In this limit a particular mixing angle follows from the quark mass dependence of the spin–orbit and tensor terms, $\phi_{HQ} = -54.7^\circ$ (35.3°), if the expectation of the heavy-quark spin–orbit interaction is positive (negative). It is often assumed that the heavy quark mixing angle holds for charmed mesons.

Fig. 2 shows the dependence of the mixing angle on the heavy quark mass for $Q\bar{u}$ and $Q\bar{s}$ mesons for the traditional and extended models. The effect of the one-loop terms is dramatic: for the $Q\bar{u}$ system the relevant spin–orbit matrix element changes sign, causing the heavy quark limit to switch from 35.3° to -54.7° . Alternatively, both models approach -54.7° in the $Q\bar{s}$ system. There is strong deviation from the heavy quark limit in both cases: $\phi(D_s) \approx \phi(D) \approx -70^\circ$. This result is not close to the heavy quark limit (which is approached very slowly)—indeed it is reasonably close to the unmixed limit of $\pm 90^\circ$!

Mixing angles can be measured with the aid of strong or radiative decays. For example, the D'_1 is a relatively narrow state, $\Gamma(D'_1) = 20.4 \pm 1.7$ MeV, while the D_1 is very broad. This phenomenon is expected in the heavy quark limit of the 3P_0 and Cornell strong decay models [2,23,24]. Unfortunately, it is difficult to exploit these widths to measure the mixing angle because strong decay models are rather imprecise.

Radiative decays are possibly more accurate probes of mixing angles because the decay vertex is established and the impulse approximation has a long history of success. Table 4 presents the results of two computations of radiative decays of D and D_s mesons. Meson wavefunctions are computed with ‘average’ parameters, as above. Transition matrix elements are evaluated in the impulse approximation and full recoil is allowed. The column labelled ‘nonrel’ reports transition matrix elements computed in the nonrelativistic limit, while the column labelled ‘rel’ contains results obtained with the full spinor structure at the photon vertex.

The nonrelativistic results can differ substantially from those of Refs. [22,23] because those computations were made in the zero recoil limit where an E1 transition, for example, is diagonal in spin. Thus the decay $D_1 \rightarrow D^*\gamma$ can only proceed via the 3P_1 component of the D_1 . Alternatively, the computations made here are at nonzero recoil and hence permit both components of the D_1 to contribute to this decay. The table entries indicate that nonzero recoil effects can be surprisingly large.

Further complicating the analysis is the large difference seen between the nonrelativistic and relativistic models (see, e.g., $D^{*+} \rightarrow \gamma D^+$). This unfortunate circumstance is due to differing signs between the heavy and light quark impulse approximation subamplitudes. Employing the full quark spinors leaves the heavy quark subamplitude largely unchanged, whereas the light quark subamplitude becomes larger, thereby reducing the full amplitude. The effect appears to be at odds with the only available experimental datum ($D^* \rightarrow D\gamma$). Clearly it would be useful to measure as many radiative transitions as possible in these sectors to better evaluate the efficacy of these (and other) models.

Once the decay model reliability has been established, ratios such as $\Gamma(D_1 \rightarrow \gamma D^*)/\Gamma(D'_1 \rightarrow \gamma D^*)$ and $\Gamma(D_1 \rightarrow \gamma D)/\Gamma(D'_1 \rightarrow \gamma D)$ will help determine the D_1 mixing angle.

4. Discussion and conclusions

A popular model of the D_s mesons is based on an effective Lagrangian description of mesonic fields in the chiral and heavy quark limits [25]. Deviations from these limits induce mass splittings which imply that the axial–vector and scalar–pseudoscalar mass differences are the same. Since the premise of this idea has been questioned in Refs. [2,26], it is of interest to consider this mass difference in the present model. Splittings for the three parameter sets considered above are shown in Table 5. Evidently, the chiral multiplet relationship holds to a very good approximation in both the D and D_s sectors and is robust against variations in the model parameters. Note that in the heavy quark limit these splittings scale as the inverse light quark mass squared with logarithmic

Table 4
Meson radiative decay rates (keV)

Mode	q_γ (MeV)	nonrel	rel	expt
$D^{*+} \rightarrow \gamma D^+$	136	1.38	0.08	1.5 ± 0.5
$D^{0*} \rightarrow \gamma D^0$	137	32.2	13.3	< 800
$D_0^+ \rightarrow \gamma D^{*+}$	361	76.0	7.55	
$D_0^0 \rightarrow \gamma D^{*0}$	326	1182	506.	
$D_1^+ \rightarrow \gamma D^{*+}$	381	$(6.34s)^2 + (3.22s + 5.9c)^2$	$(2.00s - 0.13c)^2 + (0.13s + 4.23c)^2$	
$D_1^0 \rightarrow \gamma D^{*0}$	380	$(27.05s)^2 + (19.33s + 9.63c)^2$	$(17.65s - 0.15c)^2 + (12.28s + 6.01c)^2$	
$D_1^{'+} \rightarrow \gamma D^{*+}$	381	$(6.34c)^2 + (-3.22c + 5.9s)^2$	$(2.00c + 0.13s)^2 + (-0.13c + 4.23s)^2$	
$D_1^{'0} \rightarrow \gamma D^{*0}$	384	$(27.26c)^2 + (19.35c - 9.83s)^2$	$(17.78c + 0.15s)^2 + (12.29c - 6.13s)^2$	
$D_1^+ \rightarrow \gamma D^+$	494	$(5.49s + 4.75c)^2$	$(4.17s - 0.60c)^2$	
$D_1^0 \rightarrow \gamma D^0$	493	$(8.78s + 31.42c)^2$	$(5.56s + 18.78c)^2$	
$D_1^{'+} \rightarrow \gamma D^+$	494	$(-5.49c + 4.75s)^2$	$(4.17c - 0.60s)^2$	
$D_1^{'0} \rightarrow \gamma D^0$	498	$(-8.90c + 31.41s)^2$	$(-5.62c + 18.78s)^2$	
$D_2^+ \rightarrow \gamma D^{*+}$	413	15.0	6.49	
$D_2^0 \rightarrow \gamma D^{*0}$	412	517	206	
$D_s^{*+} \rightarrow \gamma D_s^+$	139	0.20	0.00	
$D_{s0} \rightarrow \gamma D_s^{*+}$	196	6.85	0.16	
$D_{s1} \rightarrow \gamma D_s^{*+}$	322	$(1.84s)^2 + (0.99s + 2.39c)^2$	$(0.18s - 0.07c)^2 + (-0.44s + 2.13c)^2$	
$D_{s1}' \rightarrow \gamma D_s^{*+}$	388	$(2.13c)^2 + (-0.87c + 3.62s)^2$	$(0.24c - 0.10s)^2 + (0.64c + 3.19s)^2$	
$D_{s1} \rightarrow \gamma D_s^+$	441	$(2.68s + 1.37c)^2$	$(2.55s - 1.21c)^2$	
$D_{s1}' \rightarrow \gamma D_s^+$	503	$(3.54c - 1.12s)^2$	$(3.33c + 1.52s)^2$	
$D_{s2} \rightarrow \gamma D_s^{*+}$	420	1.98	3.94	

Table 5
Chiral multiplet splittings (MeV)

Params	$M(1^+(1/2^+)) - M(1^-)$	$M(0^+) - M(0^-)$
D low	411	412
D avg	391	389
D high	366	368
D_s low	384	380
D_s avg	373	370
D_s high	349	346

Table 6
Low lying bottom meson masses (MeV)

Flavour	0^-	1^-	0^+	1^+	1^+	2^+
B	5279	5322	5730	5752	5753	5759
expt	5279	5325	—	$5724 \pm 4 \pm 7$	—	5748 ± 12
B_s	5370	5416	5776	5803	5843	5852
expt	5369.6	5416.6	—	—	—	—
B_c	6286	6333	6711	6746	6781	6797
expt	6286	—	—	—	—	—

corrections due to the one-loop potentials. Thus one expects the D_s splittings to be approximately 2.8 times smaller than the D splittings. That they are only 10% smaller indicates how far these states are from the heavy quark limit.

Nevertheless, the near equivalence of these mass differences must be regarded as an accident. Indeed, the B masses given in Table 6 indicate that this relationship no longer holds. It would thus be of interest to find P-wave open bottom mesons (especially scalars). These data will distinguish chiral multiplet models from the model presented here and from more traditional constituent quark models. For example, Godfrey and Isgur claim that the B_0 meson lies between 5760 and 5800 MeV, the B_{s0} mass is 5840–5880 MeV, and the B_{c0} mass is 6730–6770 MeV. Of these, our B_{s0} mass is predicted to be 65–105 MeV lower than the Godfrey–Isgur mass.

The bottom flavoured meson spectra of Table 6 have been obtained with the ‘average’ extended model parameters and $m_b = 4.98$ GeV. As with the open charm spectra, a flavour-dependent constant was fit to each pseudoscalar. The second row reports recently measured P-wave B meson masses [27]; these are in reasonable agreement with the predictions of the first row.

When these results are (perhaps incorrectly) extrapolated to light quark masses, light scalar mesons are possible. Thus a simple $q\bar{q}$ interpretation of the enigmatic a_0 and f_0 mesons becomes feasible.

Finally, the work presented here may explain the difficulty in accurately computing the mass of the D_{s0} in lattice simulations. If the extended quark model is correct, it implies that important mass and spin-dependent interactions are present in the one-loop level one-gluon-exchange quark interaction. It is possible that current lattice computations are not sufficiently sensitive to the ultraviolet behaviour of QCD to capture this physics. The problem is exacerbated by the nearby, and presumably strongly coupled, DK continuum; which requires simulations sensitive to the infrared behaviour of QCD. Thus heavy–light mesons probe a range of QCD scales and make an ideal laboratory for improving our understanding of the strong interaction.

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